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Abstract

This paper is concerned with an adiabatic compression of a plasma sphere with a random small-scale magnetic field embedded in the plasma. The length of the field line between two intersections with the wall, determined from the random walk scaling, is large enough to make electron heat losses along the field lines negligible. Then, such a sphere may become an interesting target for Magnetized Target Fusion experiments. Key processes affecting the performance of such a target are identified and constraints on the parameters of the system are formulated.

KEYWORDS: magnetized target fusion, magneto-inertial fusion, random magnetic field, plasma liner

I. INTRODUCTION

Magnetized Target Fusion (MTF) is a version of pulsed fusion systems that relies on the slow, quasi-adiabatic compression of a magnetized plasma by a heavy liner. The plasma configurations potentially suitable for the implosions include field-reversed configurations (FRC), spheromaks, diffuse pinches, and several others (see, e. g., Ref. 1). The liner compresses the target in a shock-less fashion, with a velocity that is small compared to the sound speed in a target plasma (see Refs. [2-4], which contain also references to the earlier publications).

The range of plasma densities suitable for MTF corresponds to the densities which are much higher than those in magnetic confinement systems. In particular, the MTF plasma is typically much more collisional (see Ref. [5] for the corresponding dimensionless parameters). Still, the magnetic field plays a critically important role in suppressing heat conduction from the hot, compressed plasma to a cold liner. The liners can be made of a metal, as is the case in the on-going experiments with a Shiva-Star facility [6, 7]. They can also be made of a high-Z, cold plasma [8] (the latter approach is a variation of the plasma-liner concept described in Refs. [9] and [10]).

Creating a target with a well-defined magnetic configuration inside the liner is certainly possible, as was shown, e.g., in the Refs. [6, 7]. On the other hand, it might be beneficial for MTF if simpler targets, not requiring creation and sustainment of high-quality magnetic field in the course of slow compression, could be used. In this brief note we consider one of such possible targets: a plasma sphere with embedded random magnetic field (Fig. 1). By "random" we mean here a strongly tangled field where the field lines change direction on the scale l that is small compared to the global scale of the problem (here the sphere radius, R). This is quite different from a stochastization of the field lines in the toroidal devices caused by overlapping of resonant islands (see Sec. 1,7 in Ref. [11] and references therein), where the field line remains weakly perturbed and makes many toroidal transits before it gets displaced significantly in the radial direction. The field that we analyze here is more like the random field considered in astrophysics, in such objects as dense molecular clouds [12].

The field that we consider is weak in the sense that the ratio β of the plasma pressure to magnetic pressure is greater than one. This is a beneficial regime for plasma adiabatic compression, as the compressional *pdV* work goes then mostly to the plasma energy, with only a small fraction going to the magnetic field energy. The magnetic field, on the other hand, is assumed to be strong enough, so that plasma particles are magnetized. Then, the electron heat conduction goes predominantly along the field lines; due to the random walk of the field lines, the field line length between the core and the liner surface is much larger than the plasma radius R, and one can expect that electron heat loss to the walls will be greatly reduced.

A parameter domain where the MTF approach can be used is quite broad. In order to get specific numerical estimates, we assume that the system operates in the range of parameters discussed in [2, 3]: the initial plasma density $n_0\sim 10^{17}$-10^{19} cm $^{-3}$, the final one $n_f\sim 10^{20}$-10^{22} cm $^{-3}$, and the final plasma temperature T_f=10 keV.

The implosion can be conveniently characterized by the radial convergence C,

$$C \equiv R_0 / R_f, \tag{1}$$

with R_0 and R_f being the initial and final (at the rebound) plasma radii; the initial temperature is related to the final temperature by

$$T_0 / T_f = 1/C^2 \tag{2}$$

(see [2]). In making specific numerical estimates, we use a class of implosions with C=10, T_f=10 keV, and the fusion gain Q=10. As was shown in Ref. [13], all other parameters of the implosions can then be expressed as functions of two “input” parameters, the energy W_L deposited to the plasma by the liner, and the initial plasma radius R_0. In particular, one has (see Eqs. (19), (20) and (23) in Ref. [13]):

$$n_f\left(cm^{-3}\right) = 1000 n_0\left(cm^{-3}\right) \approx 5\times 10^{22} \frac{W_L(MJ)}{\left[R_0(cm)\right]^3}, \tag{3}$$

$$\mathrm{v}\left(cm/s\right) \approx 8\times 10^{6} \frac{W_L(MJ)}{\left[R_0(cm)\right]^2}. \tag{4}$$

This paper represents a first attempt of a scoping study of a relatively unexplored object, and one should not expect from it a comprehensive, rigorous analysis. Its result is a set of constraints that have to be imposed on the system in order it to be considered as a candidate for future, more detailed assessments.

II. COMPRESSION OF A TANGLED MAGNETIC FIELD

To obtain first rough estimates of various effects occurring during the implosion process, we use the following strategy: We assume that the thermal insulation of the imploding plasma is good, so that the compression occurs in an adiabatic regime. This allows us to obtain time histories of various parameters of the plasma and the magnetic field, by using scalings for 3D adiabatic compression [2, 13]. We then find the heat losses from this plasma, as well as the rate of the magnetic field dissipation, and thereby identify the parameter domain where the losses are indeed insignificant.

Initial spatial correlation length of the random magnetic field (Fig. 1) is assumed to be much smaller than the plasma radius, with the spatial distribution of the magnetic field being isotropic (no preferential direction for the scales significantly greater than the correlation length). As mentioned above, the magnetic field is weak (high-beta plasma) and does not affect the dynamics of the system, in particular, gives only minor contribution to the total pressure. For an homogeneous compression that the plasma experiences in our system, the distance between any two points scales as a radius R, i.e., the length l scales as:

$$\frac{l}{l_0} = \frac{R}{R_0}. \tag{5}$$

This means, in particular, that the number N of the magnetic "loops" on the scale R, remains constant in the course of the implosion,

$$N \equiv \frac{R}{l} = \frac{R_0}{l_0} = const. \tag{6}$$

The width of any magnetic flux tube scales as R, i.e., the magnetic field strength scales as $(R/R_0)^2$, and the magnetic field energy density scales as $(R/R_0)^4$. The rate of the energy density increase (by compression) is $d/dt(B^2/8\pi) = -(B^2/8\pi)4\dot{R}/R$. On the other hand, the rate of the Joule dissipation for the fields having a length-scale l is $d/dt(B^2/8\pi) = -(B^2/8\pi)/\tau_M$, where τ_M is the magnetic diffusion time over the scale l. Roughly,

$$\tau_M = \frac{l^2}{2D_M}; \; D_M = \frac{c^2}{4\pi\sigma}, \tag{7}$$

with D_M being the magnetic diffusivity and σ being the electrical conductivity. CGS system of units is used throughout this paper in the "general" equations like Eq. (), and "mixed" units are used in "practical" numerical estimates. The magnetic energy equation that accounts for both compressional build-up and resistive dissipation then becomes:

$$\frac{d}{dt}\frac{B^2}{8\pi} = \left(-\frac{4\dot{R}}{R} - \frac{1}{\tau_M}\right)\frac{B^2}{8\pi} \tag{8}$$

Note that the first term in the brackets is positive, since $\dot{R} < 0$.

As we do not want significant dissipation of the magnetic field during the implosion process, we have to impose the constraint that

$$\tau_M > R/4\dot{R}. \tag{9}$$

Specific dependence of the implosion velocity vs time is determined by the mode of the liner implosion. We will concentrate on the mode where the liner has a finite velocity at the time of its first contact with the target, and no further liner acceleration occurs. This is a mode that will be realized with the heavy plasma liner technique described in Refs. [8, 13], and in those versions of the metal liner implosions where the target is injected into the metal liner at the time when it has already reached significant velocity, similar to what is anticipated in the experiments [6, 7]. So, we assume that

$$\dot{R} \approx \mathrm{v}_0 \tag{10}$$

for the most part of the implosion, except for the very last moments near the liner rebound [13].

The length-scale of the magnetic field varies according to Eq. (5), and the magnetic diffusivity varies as $T^{-3/2} \propto R^{-2}$. Therefore, Eq (9) can be rewritten as:

$$\frac{2R_0 v_0}{D_{M0}} > N^2 \left(\frac{R}{R_0} \right)^2 . \tag{11}$$

One sees that the most severe problems with the magnetic field dissipation take place during the early stage of the implosion, when R is still not much less than R_0. This is a result of two factors: i) the temperature is low early in the implosion, meaning a low electrical conductivity; ii) with a constant implosion velocity, the compression rate, $|\dot{R}|/R$, is initially small. If condition (11) is satisfied early in the implosion, then it is satisfied until the end. Using the numerical estimate of D_{m0} from Ref. [14], and setting $R=R_0$ in Eq. (12), one finds the following numerical constraint:

$$N^2 < 5 \times 10^{-7} R_0(cm) v_0(cm/s) \left[T_0(eV) \right]^{3/2} . \tag{12}$$

For our reference case described in conjunction with Eqs. (1)-(4), one can rewrite this inequality as:

$$N^2 < 4 \times 10^3 \frac{W(MJ)}{R_0(cm)} \tag{13}$$

The resulting plots are shown by blue lines in Fig. 2.

III. ENERGY CONFINEMENT

Assuming that condition (13) is satisfied, we turn now to the issues of the heat conduction from the hot plasma to a relatively cold liner. For the magnetic field lines experiencing random changes of direction with a step size l, the field line length L between the center of the device and the wall obeys a diffusive scaling, $L \sim R^2/l \sim NR$. We assume that the electron mean free path λ is much smaller than this length (see Sec. IV). In such a case, one can evaluate the electron heat conduction time to the walls as

$$\tau_e \sim \frac{L^2}{2\chi_e} \sim N^2 \frac{R^2}{2\chi_e} , \tag{14}$$

where χ_e is the electron thermal diffusivity. It scales as $T^{5/2}/n$, so that, for the adiabatic heating ($T \propto 1/R^2$, $n \propto 1/R^3$), the r.h.s. scales as R^4. Therefore, in this case, the most severe constraints on the heat losses will appear at the last stage of the implosion, near the stagnation point.

The dwell time τ_d near the stagnation is related to the fusion gain factor Q via the Lawson criterion, $\tau_d(s) \sim 10^{14} Q / n_f\left(cm^{-3}\right)$, and the numerical estimate of χ_e reads as [13] (for the Coulomb logarithm equal to 10):

$$\chi_e(cm^2/s) \approx 6.6 \times 10^{19} \left[T(eV) \right]^{5/2} / n\left(cm^{-3}\right) . \tag{15}$$

Then the condition $\tau_e >> \tau_d$ for the final state of compression becomes:

$$N^2 >> \frac{1.3 \times 10^{44} Q}{\left[n_f\left(cm^{-3}\right) \right]^2 \left[R_f(cm) \right]^2} . \tag{16}$$

Using Eqs. (3), (4), one finds for our reference case (C=10, Q=10):

$$N^2 > \frac{50\left[R_0(cm)\right]^4}{\left[W_L(MJ)\right]^2}. \qquad (17)$$

This constraint is shown in Fig. 2 by red lines. The domain where both conditions (13) and (17) are satisfied lies between the blue and the red curves. On sees that, to have significant operational margin, it is advisable to use larger energies and smaller target radii.

Consider now particle magnetization. As we will see, in the regimes of interest, the ion gyro-radius ρ_i is smaller than the correlation length l (this meaning also that the electron gyroradius ρ_e is automatically much smaller than l). As l varies as $1/C$, and the plasma temperature and the magnetic field vary as C^2 , one sees that the ratio ρ_i/l does not depend on the convergence. So, if the condition

$$\rho_i/l<1 \qquad (18)$$

is satisfied in the initial state it will then be satisfied during the whole compression process.

The initial magnetic field strength is limited by our assumption that the magnetic pressure is smaller than the plasma pressure. In other words, we assume that

$$\beta_0 \equiv \frac{2n_0T_0}{(B_0^2/8\pi)} > 1. \qquad (19)$$

As was mentioned before, this makes the compression more efficient in that the pdV work goes mainly to the plasma heating, not to the pumping-up the magnetic field. For the gyro-radius, we will use the expression $\rho_i = \sqrt{2T/m_i}$, with m_i equal to the 2.5 proton mass. For our reference case described in the paragraphs containing Eqs. (1) - (4), the condition $\rho_i/l<1$ can be formulated as

$$N^2 < \frac{2}{\beta_0}\left(\frac{\omega_{pi0}R_0}{c}\right)^2. \qquad (20)$$

Assuming that $\beta_0=1$ and using other parameters and scalings for our reference case, one can show that this constraint can be converted to

$$N^2 < 10^5\frac{W_L(MJ)}{R_0(cm)}. \qquad (21)$$

This condition is less restrictive than Eq. (13) and is therefore subsumed by Eq. (13). In other words, in the regimes of interest for us the ions are magnetized. Their parallel heat transport is negligible compared to the electron parallel heat transport and can be neglected. With regard to the cross-field transport, for the condition $\rho_i/l<1$, it will be limited by the Bohm diffusion,

$$D_{Bohm} = \frac{1}{16}\frac{cT}{eB}. \qquad (22)$$

The role of the Bohm diffusion was assessed, in particular, in Ref. [13] and was found negligible (due to the high plasma density and small confinement times needed to reach significant Q values).

Consider now the dynamics of the alpha particles. They appear in the system near the stagnation of the liner, so we should consider their interaction with the target at its final state. Their gyroradius is approximately $\sqrt{140}$ times greater than the gyroradius of 10 keV ions with the mass of $2.5m_p$. Accordingly, the condition that the gyroradius of alphas is smaller than l in the final state is 140 times more restrictive than Eq. (21) and reads as

$$N^2 < 700\frac{W_L(MJ)}{R_0(cm)}. \tag{23}$$

The corresponding lines are shown as dashed green lines in Fig. 2. [Note that we prefer to present the results in terms of N^2 (not N), as otherwise it would be harder to see the differences in the log-log- plots of Fig. 2.]

If condition (23) is satisfied, then the lifetime of the alpha particles will be determined by their transport along the tangled field lines; it can then be evaluated as $R_f^2/l_f \mathrm{v}_\alpha = NR_0/C\mathrm{v}_\alpha$, where v_α is the alpha-particle velocity (~1.3×10^9 cm/s). One can compare it with the slowing-down time τ_{slow}, which can be presented as (see [15] and Eq. (29) in Ref. [13]): $\tau_{slow}(s) \approx 3\times10^{13}/n_f(cm^{-3}) \sim 6\times10^{-10}\left[R_0(cm)\right]^3/W_L(MJ)$. The condition that the latter is shorter than the former reads as:

$$N^2 > \frac{75\left[R_0(cm)\right]^4}{\left[W_L(MJ)\right]^2} \tag{24}$$

This condition is only slightly more restrictive than condition (17), so we will not consider it separately.

In the case where Eq. (23) is violated, evaluating the alpha life-time becomes more complex and we will not dwell on it in this first analysis. We just note that the alphas life time may still remain longer than the slowing down time. If, however, the latter condition breaks down, the system will become a batch-burn system, where alpha heating is negligible [2]. Still, with Q~10 even this mode may remain acceptable.

IV. PLASMA COLLISIONALITY

Consider an issue of the plasma collisionality. We will characterize it by the ratio of the "connection length" $R^2/l \equiv NR$ to the Coulomb mean free path λ. One can see that this dimensionless parameter decreases in the course of the compression, mostly due to the increase of the plasma temperature. So, in order to guarantee high collisionality during the course of compression, one needs to make sure that it is large in the final state. [An assumption of strong collisionality was used in the estimates of the electron cooling time above.] The condition $NR_f/\lambda_f > 1$can be presented as:

$$N^2 > 4\times10^{-3}\frac{\left[R_0(cm^{-3})\right]^2}{\left[W_L(MJ)\right]^2} \tag{25}$$

For all reasonable values of R_0 it is much weaker than condition (17) and can therefore be considered as satisfied.

One more constraint comes is related to the possible onset of anomalous resistivity. The current density j required to make magnetic field varying at a small scale l can be evaluated as

$$j \sim \frac{cB}{4\pi l}. \tag{26}$$

The relative velocity u of the electrons and ions corresponding to this current density is

$$u = \frac{j}{en} \sim \frac{cB}{4\pi enl}. \qquad (27)$$

In order to avoid the appearance of the anomalous resistivity, one has to impose a constraint (see, e.g., Sec. VII in Ref. [16] and references therein $u<\mathrm{v}_{Ti}$, where $\mathrm{v}_{Ti} = \sqrt{2T/m_i}$ is the ion thermal velocity. One can see that this condition is most difficult to satisfy early in the implosion, at $R=R_0$. The constraint can be formulated in terms of the parameter N^2 and reads as

$$N^2 < \frac{\beta_0}{2}\left(\frac{\omega_{pi0}R_0}{c}\right)^2. \qquad (28)$$

For β_0=1 It is less restrictive than condition (13) and is therefore subsumed by this condition.

V. DISCUSSION

The zero-dimensional scoping study presented in previous sections has shown that the target with a random magnetic field can be of some interest for the magnetized target fusion. It is attractive for MTF in that it eliminates any concerns related to the possible development of configurational instabilities which may affect any "regular" confinement configuration. On a macroscopic level, the compression of a sphere with small-scale random field is just a compression of a sphere of gas with a low thermal conductivity. In principle, the sphericity is not a necessary feature of this scheme: one can think of compressing prolate (or oblate) volumes of gas. This latter circumstance opens up a possibility of using shaped, magnetically driven metal liners, as in the experiments [6, 7].

Creating of an initial plasma with a small-scale, random, β~1 magnetic field immersed into it may be not a simple task. The author is not aware of any published papers where formation and characterization of such an object would be documented. An intuitively appealing way for creating such a target would be the use of numerous plasma guns generating small-scale, magnetized plasma bunches and injection of such bunches into a limited volume. This could be a version of guns envisaged in the plasma liner approach [9, 10]. Here we would need bunches with the energy per ion of only a few hundred electron-volt.

One can also try to use effects of self-generation that are present in the plasma with initially absent magnetic field but having significant temperature and density gradients. However, in this case, we will have no much control over the scale of the generated field and its magnitude.

In the MTF concept the target is supposed to be compressed by a heavy liner whose velocity is small compared to the plasma sound speed. In particular, during the first contact between the liner and the target, the liner velocity is much smaller than the sound velocity in the target. This means that the target with random fields has to be created inside an already moving liner. This imposes additional constraints on the access for the plasma bunches or plasma streams. One possibility is the use of the glide cones [13], similar to those used in or in magnetically-driven implosion of spherical liners [17] or in Fast Ignition research [18]. At any rate, creating a target inside a moving liner is a non-trivial task.

Among the physics problems not discussed in this paper is the effect of a plasma cooling near the target edge and formation of the colder transition zone between the target and the (relatively) cold liner. This problem has drawn much attention in the context of the targets with a regular magnetic field (see, e.g., Ref [5] and references therein): as the plasma beta is higher than one, a plasma pressure balance means that the colder layers near the walls must have higher density, i.e., a redistribution of the plasma over the target volume occurs in the course of the plasma compressional heating. The magnetic field is advected to this colder layer and compressed there, thereby slowing down this "cooling flow." The advection and compression of the random magnetic field may occur differently and may lead to a rapid field dissipation in the colder layer, unless the condition (13) is held by a large margin. We leave the analysis of this issue for future work.

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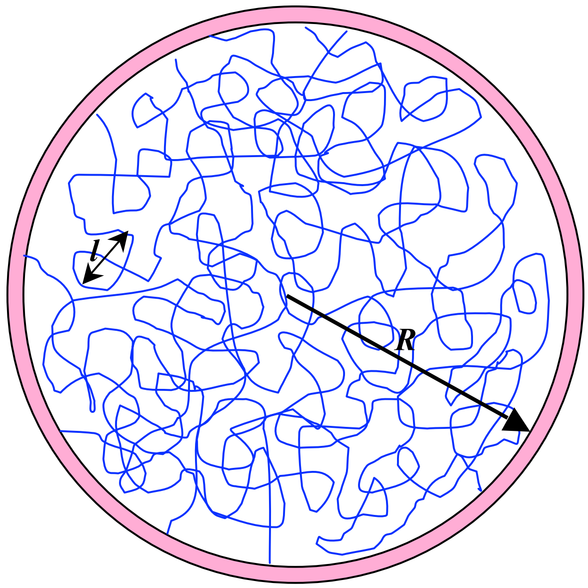

Fig. 1 Plasma sphere of a radius R with a random magnetic fieldwith a correlation length $l<<R$. Shown in blue is one of the magnetic field lines. Self-intersections are a result of the projection of a 3-dimentional object onto a plane. Thin pink shell represents a liner.

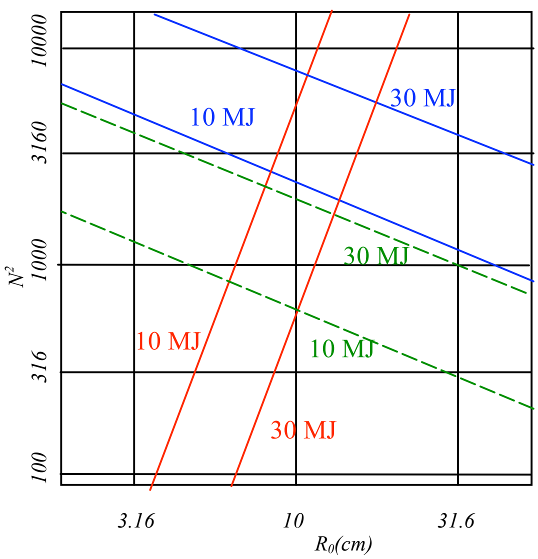

Fig. 2. The parameter space for implosions of the spherical target with random magnetic field: blue lines correspond to Eq. (13), red lines, to Eq. (17), and green lines, to Eq. (23). The numbers by the curves correspond to the liner energy, fusion yield Q is 10, radial convergence C is 10.